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Dynamical calculation of d^* mass and NN decay width in the quark delocalization, color screening model

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Abstract

The mass estimate of the $d^*(IJ^P = 03^+)$ dibaryon is improved by a dynamical calculation in the quark delocalization, color screening model. The partial decay width of d^* into an NN D-wave state is also obtained. The mass obtained is slightly larger than that obtained in adiabatic calculations, due to the anharmonicity of the effective potential between two Δ 's. The value of the width obtained due to tensor one-gluon-exchange is about 5 MeV, comparable in magnitude to earlier results found using pion exchange.

12.39.-x, 14.20.Pt, 13.75.Cs

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1 Introduction

Quantum Chromodynamics (QCD) is widely accepted as the fundamental theory of the strong interaction. However, the low energy behavior of QCD remains a challenge to understand. Lattice QCD can calculate the structure of simple quark-gluon systems but it remains extremely difficult to calculate complicated systems such as two-hadron interactions and the structure of multiquark systems. Quark models were inspirational in the development of QCD and may play a similar role in understanding the low energy behavior of QCD.

The nonrelativistic constituent quark model has been quite successful in understanding hadron spectroscopy [1]. This model has been extended by many groups to the description of baryon-baryon (B-B) interactions. There is a broad consensus that the repulsive core of the N-N interaction can be attributed to the quark internal structure of the nucleon. In this respect, the repulsive core is quite similar to that of molecular forces originating from the electronic internal structure of atoms. However, the intermediate range attraction of the N-N interaction remains absent in many quark model approaches and a scalar meson exchange has frequently been invoked to remedy this.

Spontaneous chiral symmetry breaking supports the view that Goldstone boson exchanges occur even between constituent quarks. But a question remains as to what extent this picture should be used: Should it be used for long range pion exchange only, or extended to intermediate two pion or σ exchange or even to include short range vector and other heavy meson exchange? At what point should gluon exchange between quarks be replaced? of course, a great deal of further study is needed to settle such questions definitively.

One point we would like to emphasize here is that such an approach would qualitatively differentiate between nuclear and molecular forces, which have been known for more than half a century to be similar in many features, the energy and length scale differences notwithstanding [2]. Recently, Anderson has pointed out that it is questionable to attribute the N-N interaction to meson exchange [3] except for the longest range portion of pion exchange. NN partial wave analysis can only justify long range pion exchange but not other, heavier meson exchanges [4].

In the quark model approaches mentioned above, a two-body confinement potential is usually used. We recognize that this is a highly simplified model assumption. Quark confinement is a nonperturbative property of low energy QCD. There might be important nonperturbative features missed in every two-body confinement potential model. For example, three-gluon exchange between three quarks and three-body instanton interactions do not contribute within a colorless meson or baryon, but do contribute to a multiquark system [5, 6]. There are many other types of quark interactions which result from multigluon exchanges which can not be included in a two body confinement potential. The QCD condensates within a nucleon and between two nucleons might be

different. These arguments show that a direct extension of the two body confinement to the multiquark system has not been justified even though it might be a good approximation in the single hadron case. The absence of experimentally observed color Van der Waals forces has been a problem for the two body confinement potential model for a long time. Multiquark systems provide more variations of low energy behavior of QCD and allow for further tests of the phenomenology built into quark models, especially the nature of confinement.

The quark delocalization and color screening model (QDCSM) has been developed in which an unusual confinement parameterization has been introduced. Following the concept of electron delocalization in molecular orbitals, quark delocalization is introduced to enlarge the variational Hilbert space to include various deformed six-quark configurations. The usual pair of three-quark clusters and six-quark, "bag", configurations are the two extremes of this space [7].

With only one adjustable parameter, the color screening constant μ , the existing NN ($IJ = 01, 10, 11, 00$), $N\Lambda$ ($IJ = \frac{1}{2}0, \frac{1}{2}1$) and $N\Sigma$ ($IJ = \frac{1}{2}0, \frac{1}{2}1, \frac{3}{2}0, \frac{3}{2}1$) scattering data have all been fit qualitatively [7, 8, 9, 10] within this model. A relativistic version of this model has recovered ${}^4\text{He}$, ${}^3\text{He}$, ${}^3\text{H}$ as approximate four- and three-nucleon systems starting from appropriate 12- and 9-quark configurations [11, 12, 13].

Both the relativistic and the nonrelativistic version give almost the same dibaryon spectrum, reproducing small deuteron binding and the zero energy NN $IJ = 10$ resonance. An $IJ^P = 03^+$ state, called the d^* , was found to be most interesting because of its large binding energy (~ 350 MeV) and small decay width (~ 1 MeV) [7, 14, 15, 16, 17]. This state appears in and has been studied in many models. Kamae and Fujita obtained almost the same large binding [18]. Other models obtained smaller binding [19, 20, 21].

The experimental situation regarding dibaryons remains unsettled. Precise np total cross section measurements provide a stringent limit on the NN decay width of any dibaryon in the d^* mass range [22]. Similarly, the H particle has been studied in many models and searched for experimentally for more than 20 years without any indication of its existence [23]. There is an $IJ^P = 00^-$ state, d' , which seems to have experimental support but which is hard to accommodate in a six-quark model space [24, 25, 26]. A high-mass dibaryon state predicted by Lomon *et al* [27] finds support from SATURNE pp scattering data [28].

This paper reports further study of the d^* binding and decay width with the aim of providing more reliable information for experimental searches for this dibaryon. The paper is organized as follows: a brief review of the features of the QDCSM appears in Sect. 2. Our calculation method is presented in Sect. 3. Results and discussion are provided in Sect. 4. The last section gives a summary.

2 Model Hamiltonian and wave function

The details of the QDCSM can be found in Ref.[7, 15]. Here only the model Hamiltonian and wavefunction used in our calculations are repeated.

The Hamiltonian for the 3-quark system is the same as the usual potential model [1], and for a 6-quark system, it is assumed to be

$$\begin{aligned}
H_6 &= \sum_{i=1}^6 \left(m_i + \frac{p_i^2}{2m_i} \right) - T_{CM} + \sum_{i<j=1}^6 \left(V_{ij}^C + V_{ij}^G \right), \\
V_{ij}^G &= \alpha_s \frac{\vec{\lambda}_i \cdot \vec{\lambda}_j}{4} \left[\frac{1}{r} - \frac{\pi \delta(\vec{r})}{m_i m_j} \left(1 + \frac{2}{3} \vec{\sigma}_i \cdot \vec{\sigma}_j \right) + V_T \right], \\
V_T &= \frac{1}{4m_i m_j} \left[\frac{3(\vec{\sigma}_i \cdot \vec{r})(\vec{\sigma}_j \cdot \vec{r})}{r^5} - \frac{\vec{\sigma}_i \cdot \vec{\sigma}_j}{r^3} \right], \\
V_{ij}^C &= -a_c \vec{\lambda}_i \cdot \vec{\lambda}_j \begin{cases} r^2 & \text{if } i, j \text{ occur in the same baryon orbit,} \\ \frac{1-e^{-\mu r^2}}{\mu} & \text{if } i, j \text{ occur in different baryon orbits,} \end{cases}
\end{aligned} \tag{1}$$

where $\vec{r} = \vec{r}_i - \vec{r}_j$ and all other symbols have their usual meaning except the confinement potential V_{ij}^C which will be explained below.

It is well known that confinement is a nonperturbative QCD effect; in general one does not expect it to be described by a sum of two-body interactions. We model confinement as follows: In our approach, the fractional parentage expansion method is used and the matrix elements of the six-quark Hamiltonian are simplified to a four-body overlap and a two-body matrix element. Although it cannot be demonstrated to be correct, it is at least more sophisticated than the usual, simple two-body confining interaction, and does include non-local, nonperturbative effects. We assume the following recipe to determine the two-body matrix element of confinement: The interaction takes the normal, unscreened form (quadratic in r) when the interacting quarks always remain in the same baryon orbit ϕ (see Eq.(9) below), both before and after interaction; otherwise the interaction takes the screening form (second form shown in the last of Eqs.(1)).

Even though we use a two-body interaction form to evaluate the matrix elements, our model is, nonetheless, not a potential model. It is an effective matrix element approach extended from bound states to scattering states. It reduces to the usual two-body confinement interaction within a single hadron. It does have the usual meaning of two-body interaction in the asymptotic region but not in the interacting region. The main physics introduced is the recognition that the confining interaction between two nucleons might be different from that within a nucleon. In particular, we represent the nonlocal, nonperturbative backflow of color (and pair creation, as seen in lattice studies [29]) by screening of the confining potential. If this model assumption of the confinement is not a good representation of the physics involved, we would expect that to appear in a disagreement between our calculational results and data. The quality

of the fit to B-B scattering data referred to in the introduction shows that there is no obvious contradiction. Our predictions with respect to dibaryon states will provide a further test.

We use the resonating group method (RGM) to carry out a dynamical calculation. Following the nomenclature of Ref.[25], we write the conventional ansatz for the two-cluster wavefunction as

$$|\Psi_{6q}\rangle = \mathcal{A} \sum_L \left[[\Phi_{B_1} \Psi_{B_2}]^{[\sigma]IS} \otimes \chi_L(\vec{R}) \right]^J, \quad (2)$$

where $[\sigma] = [222]$ gives the total color symmetry and all other symbols have their usual meanings. Ψ_{B_i} is the 3-quark cluster wavefunction (after removal of the center of mass motion),

$$\Psi_{B_i} = \left(\frac{2}{3\pi b^2} \right)^{3/4} \left(\frac{2}{4\pi b^2} \right)^{3/4} e^{-\left(\frac{\vec{\lambda}_i^2}{3b^2} + \frac{\vec{\rho}_i^2}{4b^2} \right)} \eta_{I_i S_i}(B_i) \chi_c(B_i), \quad (3)$$

where $\chi_c(B_i)$ is the internal color wavefunction of the baryon and the Jacobi coordinates are defined as follows,

$$\begin{aligned} \vec{\rho}_1 &= \vec{r}_1 - \vec{r}_2, & \vec{\rho}_2 &= \vec{r}_4 - \vec{r}_5, \\ \vec{\lambda}_1 &= \vec{r}_3 - \frac{1}{2}(\vec{r}_1 + \vec{r}_2), & \vec{\lambda}_2 &= \vec{r}_6 - \frac{1}{2}(\vec{r}_4 + \vec{r}_5), \\ \vec{R}_{B_1} &= \frac{1}{3}(\vec{r}_1 + \vec{r}_2 + \vec{r}_3), & \vec{R}_{B_2} &= \frac{1}{3}(\vec{r}_4 + \vec{r}_5 + \vec{r}_6), \\ \vec{R} &= \vec{R}_{B_1} - \vec{R}_{B_2}, & \vec{R}_C &= \frac{1}{2}(\vec{R}_{B_1} + \vec{R}_{B_2}). \end{aligned} \quad (4)$$

From the variational principle, after variation with respect to the relative motion wavefunction $\chi(\vec{R}) = \sum_L \chi_L(\vec{R})$, one obtains the RGM equation

$$\int H(\vec{R}, \vec{R}') \chi(\vec{R}') d\vec{R}' = E \int N(\vec{R}, \vec{R}') \chi(\vec{R}') d\vec{R}', \quad (5)$$

where $H(\vec{R}, \vec{R}')$, $N(\vec{R}, \vec{R}')$ are Hamiltonian and norm kernels, respectively. Their detailed expressions can be found in Ref.[25].

The energies, E , and the wavefunctions, $\chi(\vec{R})$, are obtained by solving the RGM equation. In practice, it is not convenient to work with the RGM expressions. We introduce generator coordinates, \vec{S}_i , to expand the relative motion wavefunction, $\chi(\vec{R})$,

$$\chi(\vec{R}) = \frac{1}{\sqrt{4\pi}} \sum_L \left(\frac{3}{2\pi b^2} \right)^{3/4} \sum_i C_{i,L} \int e^{-\frac{3}{4b^2}(\vec{R}-\vec{S}_i)^2} Y^L(\hat{\vec{S}}_i) d\Omega_{S_i}, \quad (6)$$

After the inclusion of the center of mass motion,

$$\Phi_C(\vec{R}_C) = \left(\frac{6}{\pi b^2} \right)^{3/4} e^{-\frac{3}{b^2} \vec{R}_C^2}, \quad (7)$$

the ansatz, Eq.(2), can be rewritten as

$$\begin{aligned} \Psi_{6q} = & \mathcal{A} \sum_{i,L} C_{i,L} \int \frac{d\Omega_{S_i}}{\sqrt{4\pi}} \prod_{\alpha=1}^3 \phi_{\alpha}(\vec{S}_i) \prod_{\beta=4}^6 \phi_{\beta}(-\vec{S}_i) \\ & \left[[\eta_{I_1 S_1}(B_1) \eta_{I_2 S_2}(B_2)]^{IS} Y^L(\hat{\vec{S}}_i) \right]^J [\chi_c(B_1) \chi_c(B_2)]^{[\sigma]} \end{aligned} \quad (8)$$

where $\phi_{\alpha}(\vec{S}_i), \phi_{\beta}(-\vec{S}_i)$ are the single particle orbital wavefunctions with different reference centers,

$$\begin{aligned} \phi_{\alpha}(\vec{S}_i) &= \left(\frac{1}{\pi b^2} \right)^{3/4} e^{-\frac{1}{2b^2}(\vec{r}_{\alpha} - \vec{S}_i/2)^2} \\ \phi_{\beta}(-\vec{S}_i) &= \left(\frac{1}{\pi b^2} \right)^{3/4} e^{-\frac{1}{2b^2}(\vec{r}_{\beta} + \vec{S}_i/2)^2}. \end{aligned} \quad (9)$$

With the reformulated ansatz, Eq.(8), the RGM equation Eq.(5) becomes an algebraic eigenvalue equation,

$$\sum_{j,L} C_{j,L} H_{i,j}^{L,L'} = E \sum_j C_{j,L'} N_{i,j}^{L'} \quad (10)$$

where $N_{i,j}^{L'}, H_{i,j}^{L,L'}$ are the Eq.(8) wavefunction overlaps and Hamiltonian matrix elements (without the summation over L'), respectively. By solving the generalized eigen problem, we obtain the energies of 6-quark system and corresponding wavefunctions.

In the QDCSM, the single particle orbital wavefunctions are delocalized. To implement this here, we modify Eqs.(9) as follows:

$$\begin{aligned} \phi_{\alpha}(\vec{S}_i) &\rightarrow \psi_{\alpha}(\vec{S}_i, \epsilon) = \phi_{\alpha}(\vec{S}_i) + \epsilon \phi_{\alpha}(-\vec{S}_i), \\ \phi_{\beta}(\vec{S}_i) &\rightarrow \psi_{\beta}(-\vec{S}_i, \epsilon) = \phi_{\beta}(-\vec{S}_i) + \epsilon \phi_{\beta}(\vec{S}_i), \end{aligned} \quad (11)$$

It is straightforward to extend the method to multichannel coupling. In the multichannel case, the ansatz used is

$$\begin{aligned} \Psi_{6q} = & \mathcal{A} \sum_k \sum_{i,L_k} C_{k,i,L_k} \int \frac{d\Omega_{S_i}}{\sqrt{4\pi}} \prod_{\alpha=1}^3 \psi_{\alpha}(\vec{S}_i, \epsilon) \prod_{\beta=4}^6 \psi_{\beta}(-\vec{S}_i, \epsilon) \\ & \left[[\eta_{I_1 k S_{1k}}(B_{1k}) \eta_{I_2 k S_{2k}}(B_{2k})]^{IS_k} Y^L(\hat{\vec{S}}_i) \right]^J [\chi_c(B_1) \chi_c(B_2)]^{[\sigma]} \end{aligned} \quad (12)$$

where k is the channel index. The eigen equation is similar to Eq.(10), with an additional summation over k . For example, for $IJ = 03$, we have $k = 1, 2, 3, 4$, corresponding to the channels $\Delta\Delta \ S = 3 \ L = 0$, $\Delta\Delta \ S = 3 \ L = 2$, $NN \ S = 1 \ L = 2$, and $\Delta\Delta \ S = 1 \ L = 2$.

3 Calculation method

To further simplify the calculation of the matrix elements of the six-quark Hamiltonian, the physical bases, Eq.(8), are first expanded in terms of symmetry bases (group chain classification bases). Then the powerful fractional parentage expansion method is used to calculate the matrix elements between symmetry bases of the six-quark system. Finally, the matrix elements between physical bases are obtained by the transformation between the physical bases and symmetry bases. The details can be found in Ref.[30, 31].

The partial width for d^* decay into the NN D-wave state is obtained using “Fermi’s Golden Rule”,

$$\begin{aligned}\Gamma &= \frac{1}{(2\pi)^2} \int p^2 dp d\Omega \delta(E_f - E_i) |M|^2 \\ &= \frac{1}{32\pi^2} m_{d^*} \sqrt{m_{d^*}^2 - 4m_N^2} \int |M|^2 d\Omega,\end{aligned}\quad (13)$$

where M denotes the nonrelativistic transition matrix element,

$$M = \langle d^* | H_I | \Psi_{N_1} \Psi_{N_2} e^{i\vec{p} \cdot \vec{R}} \rangle, \quad (14)$$

where \vec{R} is as above and $2p = \sqrt{m_{d^*}^2 - 4m_N^2}$ is the available relative momentum between the nucleons as determined by the energy conserving δ -function in Eq.(13). The interaction Hamiltonian, H_I , here is $H_I = V_T$.

By expanding the plane wave in terms of spherical harmonics, and taking into account angular momentum conservation, the transition matrix element can be put in the form,

$$M = -4\pi \left\langle d^* | H_I | \left[[\Psi_{N_1} \Psi_{N_2}]^{IS} Y^L(\hat{\vec{R}}) (Y^L(\hat{\vec{p}}))^* \right]^J j_L(pR + \delta_L) \right\rangle. \quad (15)$$

where we have included the phase shift, δ_L , induced by the interaction in the spherical Bessel function. Of course, the phase shift is only apparent at large separations.

On the other hand, the relative motion wavefunction between two nucleons can be determined from either the RGM equation or the algebraic equation, Eq.(10). The relative motion wavefunction can be written as

$$\chi_L(\vec{R}) = \sqrt{4\pi} \sum_i C_i \left(\frac{3}{4\pi b^2} \right)^{3/4} e^{-\frac{3}{4b^2}(\vec{R}^2 + \vec{S}_i^2)} i_L\left(\frac{3}{2b^2} R S_i\right) Y^L(\hat{\vec{R}}), \quad (16)$$

where $i_L(z)$ is the modified spherical Bessel function. If we can match the expansion

$$\sum_i C_i \left(\frac{3}{4\pi b^2} \right)^{3/4} e^{-\frac{3}{4b^2}(\vec{R}^2 + \vec{S}_i^2)} i_L\left(\frac{3}{2b^2} R S_i\right)$$

to the function $j_L(pR + \delta_L)$ for the separation in the calculated region, then the wavefunction used in Eq.(15) is just the wavefunction defined in Eq.(8) with an additional factor of $\sqrt{4\pi}(Y^L(\hat{\vec{p}}))^*$, which contributes a factor 4π to the decay width after the integration over the angular space, Ω .

We have used a type of box normalization method to calculate the transition matrix element. Because the wavefunction of the initial bound state d^* is compact, its amplitude at large separation ($R > R_0 \sim 3.3$ fm) can be safely neglected. By contrast, the final NN D-wave state with the same energy as that of the d^* is a scattering state, so it extends over the entire space. However, the transition matrix element depends only on the overlap between the d^* and NN D-wave states, so only that part of the wavefunction of the NN D-wave with separation less than R_0 fm can contribute. To determine the wave function of the final state NN D-wave in this finite range, the RGM equation is used so that the final state interaction is taken into account automatically. By adjusting the box size, the correct internal relative momentum NN D-wave state could be obtained to meet the requirement of energy conservation in the decay. Then a normalization correction due to the finite range of the wavefunction is needed; this is obtained from the RGM calculation. In this way, we could tackle the bound state problem and the scattering state problem on the same footing. For example, the phase shifts of the scattering state can be obtained from the comparison of the calculated wavefunction with the free Bessel function, $j_L(pR)$.

4 Results and discussion

As a test of the model, we first attempted a calculation of the deuteron. The parameters we used in the calculation, which are fixed by baryon properties and NN scattering, are

$$m_u = m_d = 313 \text{ MeV}, \quad b = 0.603 \text{ fm}, \quad a = 25.13 \text{ MeV/fm}^2, \quad \alpha_s = 1.54, \quad \mu = 1.0 \text{ fm}^{-2},$$

These parameter values are very similar to those usually appearing in constituent quark models [1]. The single channel and channel coupling calculations do show an attraction between nucleons, but it is insufficient to form a bound state. The mixing of the D-wave into the deuteron state is also too small (less than 1%) to affect the result. Together with the absence of a long range tail for the NN interaction, this serves to emphasize the need for the long range part of one pion exchange (OPE), which is absent from our model. To reproduce the deuteron, which is an extremely extended object, it will be necessary to extend the QDCSM to include OPE. To avoid double counting, a short-distance cutoff of the pion-quark coupling is also required. We will address these issues in a future publication [33].

The case of the d^* is quite different, however. The small size and the large delocalization of quarks shows that it is a highly compact object. The effect of OPE should be

considerably reduced compared to the case of the deuteron. Using the same parameters, a dynamical calculation of the d^* was carried out. These results appear in Table I. We find that the mass of d^* is somewhat larger than that obtained by our earlier adiabatic calculation [15], mainly because the relative motion energy found is larger than that of the zero-point oscillation energy used in the adiabatic calculation. Although this d^* mass exceeds the $NN\pi\pi$ threshold, it should remain a narrow resonance due to the extremely small phase space available [32].

Table I. Mass of the d^* in the QDCSM.
sc= single channel; cc2 = $\Delta\Delta$ S-wave + NN D-wave; cc4 = full four-channel coupling.

	sc	cc1	cc2	adiabatic [15]
mass (MeV)	2186	2180	2176	2134
$\sqrt{\langle r^2 \rangle}$ (fm)	1.2	1.2	1.3	1.4

We have also calculated the partial decay width of d^* into the NN D-wave, using the method introduced in last section. The result is dominated by the stretched ($m_J = \pm 3$) states which contribute $\Gamma^{\text{str.}} = 1.95$ MeV each to the sum over spin projections. We find $\Gamma = 4.32$ MeV for the total width. This result is similar to that obtained from meson exchange [17]. It should be expected that the total width will be increased by including OPE as discussed above for the deuteron. Although it is generally believed that the tensor interaction due to OPE is much stronger than that from one gluon exchange, we expect this increase will be limited due to the combination of the compactness of the d^* and the required short-distance cutoff of the pion-quark coupling.

5 Summary

In the framework of QDCSM, dynamical calculations of the d^* dibaryon and the partial decay width of the d^* into an NN D-wave state were carried out. The results include a slight increase in the mass of the d^* over the adiabatic result, rising above the $NN\pi\pi$ threshold. However, the small phase space still suggests that the resonance will be narrow and, indeed, the calculated $d^* \rightarrow NN$ D-wave decay width is on the order of a few MeV.

Our calculation of the deuteron shows the predictive power of the QDCSM on the one hand, but on the other, that OPE must be included in the calculation of extended objects (with small delocalization). However, we do not expect that the addition of OPE will have a large effect on the calculation of d^* , because of its compactness. Addition of OPE to the QDCSM is underway and preliminary results show that both the D-wave component and the radius of deuteron are indeed improved. Details will be presented in a forthcoming paper [33].

The calculated $d^* \rightarrow NN$ D-wave decay width is larger than allowed by the Lisowski [22] data. However for a state as high as 2.2 GeV, sea quark excitation or the $NN\pi$ component may well be important. Inclusion of this component may affect the mass and decay width strongly. In particular, the $d^* \rightarrow NN\pi$ decay width may be significantly increased [32]. We believe that it is still too early to conclude that the d^* has been ruled out by precise np total cross section measurements and that further calculations including these complications are needed.

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References

- [1] N. Isgur and G. Karl, Phys. Lett. **72B**, 109 (1977); **74B**, 353 (1978); Phys. Rev. **D18**, 4187 (1978); **D19**, 2653 (1979); **D20**, 768 (1979); **D20**, 1191 (1979); **D21**, 1868 (1980).
- [2] A. Bohr and B.R. Mottelson, Nuclear Structure, Vol.I, (W.A. Benjamin, Inc., New York, 1965), p.269.
- [3] P. W. Anderson, Phys. Today, **53**, 11 (2000).
- [4] R. Timmermans, in *Properties and Interactions of hyperons*, ed. B. F. Gibson *et al*, (World Scientific, Singapore, 1994) p.179.
- [5] W. Lucha, F. F. Schöberl and D. Gromes, Phys. Rep. **200**, 127 (1991).
- [6] B. Ch. Metsch, Baryons'98, eds. D. W. Menze and B. Ch. Metsch (World Scientific, Singapore, 1999) p.63-64.
- [7] F. Wang, G. H. Wu, L. J. Teng and T. Goldman, Phys. Rev. Lett. **69**, 2901 (1992).
- [8] G. H. Wu, L. J. Teng, J. L. Ping, F. Wang and T. Goldman, Phys. Rev. **C53**, 1161 (1996).
- [9] G. H. Wu, L. J. Teng, J. L. Ping, F. Wang and T. Goldman, Modern Phys. Lett., **A10**, 1895 (1995).
- [10] G. H. Wu, J. L. Ping, L. J. Teng, F. Wang and T. Goldman, Nucl. Phys. **A673**, 279 (2000).
- [11] T. Goldman, K. Maltman, G.J. Stephenson, Jr. and K.E. Schmidt, Nucl. Phys. **A481**, 621 (1988).

- [12] T. Goldman, in "Nuclear Chromodynamics", eds. S. Brodsky and E. Moniz, (World Scientific, Singapore, 1986), p.363.
- [13] K. Maltman, G.J. Stephenson, Jr. and T. Goldman, Phys. Lett. **B324**, 1 (1994).
- [14] T. Goldman, K. Maltman, G. J. Stephenson, Jr., K. E. Schmidt and F. Wang, Phys. Rev. **C39**, 1889 (1989).
- [15] F. Wang, J. L. Ping, G. H. Wu, L. J. Teng and T. Goldman, Phys. Rev. **C51**, 3411 (1995).
- [16] T. Goldman, K. Maltman, G. J. Stephenson, Jr., J. L. Ping and F. Wang, Mod. Phys. Lett. **A13**, 59 (1998).
- [17] C. W. Wong, Phys. Rev. **C57**, 1962 (1998).
- [18] Y. Kamae and T. Fujita, Phys. Rev. Lett. **38**, 471 (1977).
- [19] P. J. G. Mulders, *et al.*, Phys. Rev. Lett. **40**, 1543 (1978); Phys. Rev. **D21**, 2653 (1980).
- [20] K. Yazaki, Prog. Theor. Phys. Suppl. **91**, 146 (1987).
- [21] N. R. Walet, Phys. Rev. **C48**, 2222 (1993).
- [22] P. Lisowski et al., Phys. Rev. Lett. **49**, 255 (1982).
- [23] R. L. Jaffe, Phys. Rev. Lett. **38**, 195 (1977).
- [24] R. Bilger, H. A. Clement and M. G. Schepkin, Phys. Rev. Lett. **71**, 42 (1993); **72**, 2792 (1994).
- [25] G. Wagner, L. Y. Glozman, A. J. Buchmann and A. Faessler, Nucl. Phys. **A594**, 263 (1995).
- [26] J. L. Ping, F. Wang and T. Goldman, *d' in the quark delocalized, color screening model*, submitted to Phys. Rev. **C**.
- [27] P. Gonzalez, P. LaFrance and E. L. Lomon, Phys. Rev. **D35**, 2142 (1987).
- [28] F. Lehar, Baryons'98, eds. D. W. Menze and B. Ch. Metsch (World Scientific, Singapore, 1999) p.622.
- [29] B. Svetitsky, in *Nuclear Chromodynamics*, eds. S. Brodsky and E. Moniz (World Scientific), Singapore (1986) pp. 80-81; E. Laermann, F. Langhammer, I. Schmitt, and P. M. Zerwas, *ibid.*, pp. 117-133; Phys. Lett. **B173**, 437 (1986); K. D. Born, et al., Phys. Rev. **D40**, 1653 (1989).

- [30] J. Q. Chen *et al.* Nucl. Phys. **A393**, 122 (1983).
- [31] F. Wang, J. L. Ping and T. Goldman, Phys. Rev. **C51**, 1648 (1995).
- [32] C. W. Wong, Phys. Rev. **C58**, 2414 (1998).
- [33] J. L. Ping, F. Wang and T. Goldman, in preparation.